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Long-range Order in the A-like Phase of Superfluid ^3He in Aerogel

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Abstract A mutual action of the random anisotropy brought in the superfluid ^3He by aerogel and of the global anisotropy caused by its deformation is considered. Strong global anisotropy tends to suppress fluctuations of orientation of the order parameter and stabilizes ABM order parameter. In a limit of vanishing anisotropy fluctuations of ABM order parameter became critical. It is argued that still in a region of small fluctuations the order parameter changes its form to be less sensitive to the random anisotropy. For a favorable landscape of the free energy of superfluid ^3He the fluctuations remain small even in a limit of vanishing global anisotropy and the long-range order is maintained.

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1 Introduction

Recent NMR experiments with the superfluid ^3He in a uniaxially compressed aerogel¹ have shown that the state of the A-like phase is very sensitive to a global anisotropy of aerogel induced by its deformation. The global anisotropy stabilizes long-range order in a contrast to the random local anisotropy which tends to disrupt this order. The mechanism of disruption of a long-range order is the unlimited growth of fluctuations of the order parameter in directions of its degeneracy (Goldstone fluctuations)^{2,3}. In the case of superfluid ^3He these are fluctuations of orientation of the order parameter. Deformation of aerogel gives rise to the global anisotropy which lifts degeneracy of the order parameter of superfluid ^3He with

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respect to the orbital rotations. The lifting of degeneracy tends to suppress the Goldstone fluctuations. So, the state of superfluid ^3He in a deformed aerogel is a result of competition between the random local and regular global anisotropy. In a limit of strong global anisotropy the Goldstone fluctuations are small and the order parameter of the A-like phase has ABM form^{1,4}. In a limit of vanishing anisotropy, if the form of the order parameter is fixed and only its orientation can vary, a possible result of disruption of orientational long-range order is transition in the Larkin-Imry-Ma (LIM) state⁵. A straightforward interpolation between the two limits does not exhaust possibilities of variation of a state of ^3He in aerogel with a change of global anisotropy. There exist a feedback effect of fluctuations on a form of the order parameter. Depending on a landscape of the free energy of superfluid ^3He in a vicinity of the ABM order parameter this effect can be significant. Variation of a form of the order parameter of ^3He -A under the influence of fluctuations adds another dimension to the manifold of possible states of this phase. This possibility was considered previously only for the isotropic aerogel⁶. In the present paper the global anisotropy is introduced in this analysis as an external parameter. It is shown in particular, that if both anisotropy of aerogel and variation of a form of the order parameter are taken into account, the long-range order in the A-like phase of superfluid ^3He can be preserved even in a limit of a vanishing global anisotropy.

2 Effect of anisotropy

Interaction of aerogel with the superfluid ^3He is described phenomenologically by the extra term in the Ginzburg and Landau functional:

$$F_\eta = N(0) \int \eta_{jl}(\mathbf{r}) A_{\mu j} A_{\mu l}^* d^3 r, \quad (1)$$

where $N(0)$ is the density of states at the Fermi level, $A_{\mu j}$ – the order parameter and $\eta_{jl}(\mathbf{r})$ – the random anisotropy tensor. On the strength of $t \rightarrow -t$ invariance tensor $\eta_{jl}(\mathbf{r})$ is real and symmetric. For isotropic aerogel the average $\langle \eta_{jl}(\mathbf{r}) \rangle = 0$. To account for a possible global anisotropy of aerogel a constant (\mathbf{r} -independent) symmetric tensor κ_{jl} has to be added to $\eta_{jl}(\mathbf{r})$. The resulting expression for the GL free energy has the following structure:

$$F_{GL} = N(0) \int d^3 r [f_0 + f_V + (\eta_{jl}(\mathbf{r}) + \kappa_{jl}) A_{\mu j} A_{\mu l}^*]. \quad (2)$$

Here

$$f_0 = \tau A_{\mu j} A_{\mu j}^* + \frac{1}{2} \sum_{s=1}^5 \beta_s I_s \quad (3)$$

is the unperturbed, or “bare” GL free energy, I_s - 4-th order invariants in the expansion of the free energy over $A_{\mu j}$. Coefficients β_1, \dots, β_5 are phenomenological constants⁷. Tensors $\eta_{jl}(\mathbf{r})$ and κ_{jl} can be defined as traceless, i.e. their traces are included in the definition of $\tau = (T - T_c)/T_c$.

For the gradient energy f_V we take a model isotropic expression

$$f_V = \frac{2\xi_0^2}{5} \left(\frac{\partial A_{\mu l}}{\partial x_n} \frac{\partial A_{\mu l}^*}{\partial x_n} \right), \quad (4)$$

where $\xi_0 = \hbar v_F / (2\pi T_c)$ is the coherence length in the superfluid state. Equilibrium configuration of the order parameter is found from the equation

$$\frac{\partial f_0}{\partial A_{\mu j}^*} - \frac{2\xi_0^2}{5} \left(\frac{\partial^2 A_{\mu j}}{\partial x_n^2} \right) + \kappa_{lj} A_{\mu l} = -A_{\mu l} \eta_{lj}(\mathbf{r}) \quad (5)$$

and its complex conjugated. For high porosity aerogel tensor $\eta_{lj}(\mathbf{r})$ can be treated as a small perturbation. Solution of Eq.(5) can be sought as a sum of the average order parameter $\bar{A}_{\mu j}$ and of a small fluctuation $a_{\mu j}(\mathbf{r})$:

$$A_{\mu j} = \bar{A}_{\mu j} + a_{\mu j}(\mathbf{r}). \quad (6)$$

$\bar{A}_{\mu j}$ is assumed to be not far from one of the minima of f_0 . The long-range order exist when the average order parameter is finite.

Following the standard perturbation procedure⁸ we expand Eq. (5) up to the second order in $a_{\mu j}(\mathbf{r})$ and $\eta_{jl}(\mathbf{r})$. The linear terms render equations for the fluctuations:

$$\frac{\partial^2 f_0}{\partial A_{\mu j}^* \partial A_{vl}} a_{vl} + \frac{\partial^2 f_0}{\partial A_{\mu j}^* \partial A_{vl}^*} a_{vl}^* - \frac{2\xi_0^2}{5} \left(\frac{\partial^2 a_{\mu j}}{\partial x_n^2} \right) + \kappa_{lj} a_{\mu l} = -\eta_{lj} \bar{A}_{\mu l}, \quad (7)$$

$$\frac{\partial^2 f_0}{\partial A_{\mu j} \partial A_{vl}^*} a_{vl}^* + \frac{\partial^2 f_0}{\partial A_{\mu j} \partial A_{vl}} a_{vl} - \frac{2\xi_0^2}{5} \left(\frac{\partial^2 a_{\mu j}^*}{\partial x_n^2} \right) + \kappa_{lj} a_{\mu l}^* = -\eta_{lj} \bar{A}_{\mu l}^*, \quad (8)$$

and the average of Eq. (5) over the ensemble of $\eta_{jl}(\mathbf{r})$ – the equation for the $\bar{A}_{\mu j}$:

$$\begin{aligned} \frac{\partial f_0}{\partial A_{\mu j}^*} + \frac{1}{2} \left[\frac{\partial^3 f_0}{\partial A_{\mu j}^* \partial A_{vl} \partial A_{\beta m}} \langle a_{vl} a_{\beta m} \rangle + 2 \frac{\partial^3 f_0}{\partial A_{\mu j}^* \partial A_{vl} \partial A_{\beta m}^*} \langle a_{vl} a_{\beta m}^* \rangle \right] + \\ \langle \eta_{jl} a_{\mu l} \rangle + \kappa_{lj} \bar{A}_{\mu l} = 0. \end{aligned} \quad (9)$$

The average $\langle \eta_{jl} a_{\mu l} \rangle$ can be combined with $\tau \bar{A}_{\mu j}$ in $\frac{\partial f_0}{\partial A_{\mu j}^*}$. The remaining averages of binary products of fluctuations i.e. $\langle a_{vl} a_{\beta m} \rangle = \langle a_{vl}(\mathbf{r}) a_{\beta m}(\mathbf{r}) \rangle$ yield corrections to the order parameter.

The state of the unperturbed superfluid ^3He is continuously degenerate with respect to separate rotations in spin and in orbital spaces. The latter is of significance here. The random anisotropy $\eta_{jl}(\mathbf{r})$ breaks locally rotational degeneracy and induces fluctuations $a_{\mu j}(\mathbf{r})$. The “longitudinal” fluctuations, which change the magnitude and the form of the order parameter are weakly effected by the global anisotropy. Their binary averages were estimated before⁹

$$\langle a_{\mu j} a_{vn} \rangle \sim \frac{1}{8\pi} \frac{\Phi_{jlmn}(0)}{\xi_0^3} \frac{\bar{A}_{\mu l} \bar{A}_{vm}}{\sqrt{2|\tau|}}. \quad (10)$$

Here

$$\Phi_{jlmn}(0) = \left[\int \langle \eta_{jl}(\mathbf{k}) \eta_{mn}(-\mathbf{k}) \rangle \frac{do}{4\pi} \right]_{k=0} = \Phi_0 (\delta_{jm} \delta_{ln} + \delta_{jn} \delta_{lm} - \frac{2}{3} \delta_{jl} \delta_{mn}).$$

Integral in the square brackets is taken over the solid angle do in \mathbf{k} -space. The relative value of these fluctuations with respect to the square of the average order parameter is characterized by the parameter $g_\tau = \Phi_0 / (\xi_0^3 \sqrt{|\tau|})$. For aerogel with the radius of strands ρ and the average distance between them ξ_a $g_\tau \sim \rho^2 / (\xi_0 \xi_a \sqrt{|\tau|})$, which is small if the temperature T is not too close to T_c .

Effect of fluctuations of orientation of the order parameter, or transverse fluctuations does depend on a global anisotropy. Let us start with a “strongly” compressed aerogel when definitely $\bar{A}_{\mu j} = A_{\mu j}^{ABM}$:

$$A_{\mu j}^{ABM} = \Delta \frac{1}{\sqrt{2}} \hat{d}_\mu (\hat{m}_j + i \hat{n}_j). \quad (11)$$

Here d_μ is a unit vector in spin space, \mathbf{m} and \mathbf{n} - two mutually orthogonal unit vectors in orbital space. In a uniaxially compressed aerogel vector $\mathbf{l} = \mathbf{m} \times \mathbf{n}$ is oriented along the direction of compression, which will be taken as z -axis. Then tensor κ_{jl} is diagonal, with the components $\kappa_{xx} = \kappa_{yy} = -\kappa$, $\kappa_{zz} = 2\kappa$, $\kappa > 0$. To obtain equation for the transverse fluctuations we have to multiply Eq.(7) by $\frac{\partial \bar{A}_{\mu j}^*}{\partial \theta_q} = e^{jqn} \bar{A}_{\mu n}^*$, where e^{jqn} is antisymmetric tensor, Eq.(8) by $\frac{\partial \bar{A}_{\mu j}}{\partial \theta_q} = e^{jqn} \bar{A}_{\mu n}$ and to sum the obtained equations. Vector θ_q specifies infinitesimal rotation of the order parameter. The resulting equation is

$$\begin{aligned} \frac{\partial \bar{A}_{\mu j}^*}{\partial \theta_q} \kappa_{jl} a_{\mu l} + \frac{\partial \bar{A}_{\mu j}}{\partial \theta_q} \kappa_{jl} a_{\mu l}^* - \frac{2\xi_0^2}{5} \frac{\partial^2}{\partial x_n^2} \left(\frac{\partial \bar{A}_{\mu j}^*}{\partial \theta_q} a_{\mu j} + \frac{\partial \bar{A}_{\mu j}}{\partial \theta_q} a_{\mu j}^* \right) = \\ - \frac{1}{2} \eta_{jl} \frac{\partial}{\partial \theta_q} (\bar{A}_{\mu j}^* \bar{A}_{\mu l} + \bar{A}_{\mu j} \bar{A}_{\mu l}^*). \end{aligned} \quad (12)$$

Combinations $\frac{\partial \bar{A}_{\mu j}^*}{\partial \theta_q} a_{\mu j} + \frac{\partial \bar{A}_{\mu j}}{\partial \theta_q} a_{\mu j}^*$ are transverse fluctuations.

Using $\bar{A}_{\mu j}$ given by Eq.(11) and taking Fourier transform of $a_j(\mathbf{r}) = d_\mu a_{\mu j}(\mathbf{r})$ we arrive at the following expression for the only finite transverse component $a_j(\mathbf{k})$:

$$l_j a_j(\mathbf{k}) = - \frac{5\sqrt{2}\Delta}{4(5\kappa + \xi_0^2 k^2)} [l_j \eta_{jl}(\mathbf{k})(m_l + i n_l)] \quad (13)$$

The only non-vanishing average in Eq. (9) originating from the transverse fluctuations is:

$$\langle a_3(0) a_3^*(0) \rangle = \frac{25}{8} \int \frac{\Delta^2 \Phi_0}{(5\kappa + \xi_0^2 k^2)^2} \frac{k^2 dk}{\pi^2} = \frac{5\sqrt{5}\Delta^2 \Phi_0}{32\pi \xi_0^3 \sqrt{\kappa}}. \quad (14)$$

The disorder can be treated as a perturbation when the fluctuation is small, i.e.

$$\frac{\langle a_3(0) a_3^*(0) \rangle}{\Delta^2} \equiv g_\kappa \ll 1. \quad (15)$$

With the decreasing κ parameter $g_\kappa = \frac{5\sqrt{5}\phi_0}{32\pi\xi_0^3\sqrt{\kappa}}$ grows as $1/\sqrt{\kappa}$. Perturbation theory approach breaks down at $g_\kappa \sim 1$. At smaller anisotropy transverse motion of the order parameter can not be described within the mean field approach. Situation is analogous to the critical region in a vicinity of a temperature of a continuous phase transition, except that in the case of a weak quenched disorder only transverse fluctuations are critical. Longitudinal fluctuations remain small and a short-range order can be preserved. Intensity of fluctuations is controlled by the global anisotropy κ , which in the present case is analogous to parameter $\tau = (T - T_c)/T_c$ for thermal fluctuations. The condition $g_\kappa \sim 1$ can be used for an order of magnitude estimation of a borderline anisotropy κ_c below which transverse fluctuations became critical. Considering aerogel as a collection of randomly distributed pieces of strand of a length ε and of a radius ρ with the average porosity P and using results of the Rainer and Vuorio theory of “small objects” in superfluid ^3He one can obtain the following estimations¹⁰: $\Phi_0 \sim \varepsilon\xi_0^2(1-P)$, $\kappa \sim \gamma(1-P)(\xi_0/\rho)$. Transverse fluctuations are critical if deformation $\gamma < \gamma_c \equiv \rho\varepsilon^2(1-P)/\xi_0^3$. For comparison with the Ref.⁵ let us substitute $\varepsilon = \xi_a$ as it is assumed there. Here ξ_a is the average distance between the strands, introduced as $\pi\rho^2/\xi_a^2 = (1-P)$. With this assumption $\gamma_c \sim (\rho/\xi_0)^3$. When expressed in terms of ξ_a and Larkin-Imry-Ma length L_{LIM} the borderline deformation $\gamma_c \sim (\xi_a/L_{LIM})^{3/2}$ coincides with the deformation at which transition from the uniform ABM state to the LIM state is predicted in Ref.⁵. It means that the predicted transition falls into the region where transverse fluctuations are critical. The mean-field picture used for the prediction of the transition does not apply in this region and can be used only as a qualitative guidance. An adequate description of a possible transition and of the emerging state have to be based on the formalism used for description of critical phenomena. Renormalization group analysis of several other systems with a quenched random anisotropy, in which formation of LIM state would be expected on a basis of the mean-field argument, proves that a state with the quasi long-range order (QLRO) forms instead.¹¹ In the QLRO state the average order parameter is zero, but decay of local correlations of the order parameter with a distance obeys a power law as it is expected for a decay of correlations in a critical point.

The order of magnitude estimation of the borderline deformation for $\rho/\xi_0 \sim (1/10)$ yields $\gamma_c \sim 10^{-3}$ as in Ref.⁵. Quantitative treatment¹⁰ of the model of strands within the Rainer and Vuorio theory brings this estimation down to $\gamma_c \sim 10^{-4} \div 10^{-5}$, i.e. a very high level of isotropy is required for observation of critical phenomena in the considered system. But, as it was pointed out before⁹ a deviation of the order parameter of the A-like phase from the ABM form can start in a region where transverse fluctuations of the order parameter are still small and the perturbation theory does apply.

3 Effect of fluctuations

For anisotropy κ within the interval $\kappa_c \ll \kappa \ll \tau$ transverse fluctuations are small but still much greater then the longitudinal: $g_\tau \ll g_\kappa \ll 1$. The estimated critical anisotropy $\kappa_c \sim 10^{-5} \div 10^{-6}$ and $\tau \sim 0.1$, so the interval is wide. Within this interval contribution of the longitudinal fluctuations to Eq. (9) can be neglected. That simplifies calculation of corrections to the order parameter. Substitution of

expression (14) for fluctuations and Eq. (11) as the average order parameter in Eq. (9) renders an equation for the gap Δ of the ABM phase corrected for the transverse fluctuations:

$$\tau + \beta_{245}(1 + g_\kappa)\Delta^2 = 0. \quad (16)$$

It differs from the analogous equation for the unperturbed ABM-phase by the extra factor $(1 + g_\kappa)$ in front of a sum of the coefficients $\beta_{245} = \beta_2 + \beta_4 + \beta_5$. Parameter g_κ is positive by its definition. Fluctuations depress Δ^2 and the condensation energy of the ABM-phase in comparison with the unperturbed case by a factor $(1 + g_\kappa)^{-1}$, i.e. the renormalized condensation energy $f(A_{\mu j}^{ABM}) = f_0(A_{\mu j}^{ABM})/(1 + g_\kappa)$.

The amount for which the condensation energy is depressed depends on a coupling of the average order parameter to the random anisotropy. There exist a class of orbitally isotropic, or “robust” order parameters for which the random anisotropy does not excite transverse fluctuations and there is no ensuing suppression of their condensation energy ($g_\kappa = 0$). That happens when the driving term in the r.h.s. of Eq. (12) vanishes:

$$\frac{d}{d\theta_q} (A_{\mu j} A_{\mu l}^* + A_{\mu l} A_{\mu j}^*) = 0. \quad (17)$$

This condition means that the combination in the brackets does not change at an arbitrary infinitesimal rotation θ_q , i.e. this combination is proportional to the unit tensor:

$$A_{\mu j} A_{\mu l}^* + A_{\mu l} A_{\mu j}^* \sim \delta_{jl}. \quad (18)$$

An immediate example of the robust order parameter is that of BW. Transverse fluctuations favor robust order parameters over non-robust. One can conclude that when the global anisotropy is weak the transverse fluctuations induced by aerogel tend to favor BW phase over the ABM and to shrink a region of stability of the ABM phase in comparison with the bulk liquid.

Returning to the A-like phase we have to take into account that it is an equal spin pairing state. Among these states the one satisfying condition (17) up to an arbitrary rotations in spin and in orbital spaces corresponds to the A-like robust order parameter⁹:

$$A_{\mu j}^R = \Delta \frac{1}{\sqrt{3}} [\hat{d}_\mu (m_j + i n_j) + \hat{e}_\mu l_j], \quad (19)$$

where $\mathbf{m}, \mathbf{n}, \mathbf{l}$ are mutually orthogonal orbital unit vectors, \mathbf{d}, \mathbf{e} – mutually orthogonal unit spin vectors. This order parameter is not a minimum of the “bare” free energy f_0 . The relative difference of “bare” energies of the robust and ABM-states $\varepsilon_0 \equiv [f_0(A_{\mu j}^R) - f_0(A_{\mu j}^{ABM})]/f_0(A_{\mu j}^{ABM})$ can be expressed in terms of the coefficients β_1, \dots, β_5 : $\varepsilon_0 = (\beta_{13} - 4\beta_{45})/(9\beta_2 + \beta_{13} + 5\beta_{45})$. For the weak coupling values of β -coefficients this ratio is 1/19, i.e. the density of the “bare” free energy of the robust state is only slightly higher than that of the ABM-state. Assume that the strong coupling corrections to β_1, \dots, β_5 leave ε_0 small. The relative difference of renormalized energies of the two states $\varepsilon \equiv [f(A_{\mu j}^R) - f(A_{\mu j}^{ABM})]/f(A_{\mu j}^{ABM})$ depends on the global anisotropy κ via parameter g_κ : $\varepsilon = \varepsilon_0 - g_\kappa + \varepsilon_0 g_\kappa$. According to Eq. (14) $g_\kappa \sim 1/\sqrt{\kappa}$. At sufficiently small κ when $g_\kappa > \varepsilon_0/(1 - \varepsilon_0)$, $\varepsilon < 0$ and the robust state became energetically more favorable than the ABM. That happens at

$g_\kappa \approx \varepsilon_0 \ll 1$, i.e. the transverse fluctuations are still small and the perturbation theory does apply. In terms of a global anisotropy condition $g_\kappa \approx \varepsilon_0$ corresponds to $\kappa \approx \kappa_c/\varepsilon_0^2 \gg \kappa_c$. Comparison of free energies indicates a possibility of a discontinuous transition from the ABM into the robust state or in a state with even lower free energy when the global anisotropy decreases. A landscape of the free energy of superfluid ^3He is not yet established. That impedes a definitive prediction of a character and position of transition in the robust state. Continuous change of a form of the order parameter as a function of anisotropy can not be excluded too. As an illustration of possible changes of a form of the order parameter consider an interpolation between the ABM and the robust order parameters:

$$A_{\mu j}^{int} = \frac{\Delta}{\sqrt{3+2v^2}} [(1-iv)d_\mu(m_j + in_j) + e_\mu l_j]. \quad (20)$$

At $v \rightarrow \infty$ $A_{\mu j}^{int}$ goes over into $A_{\mu j}^{ABM}$ and at $v = 0$ – into $A_{\mu j}^R$. Coefficient v is a “fraction” of the ABM-order parameter in $A_{\mu j}^{int}$. Coupling of the $A_{\mu j}^{int}$ with global anisotropy is determined by a combination

$$A_{\mu j}^{int}(A_{\mu l}^{int})^* \kappa_{jl} = -\frac{\Delta^2}{3+2v^2} v^2 l_j l_l \kappa_{jl}. \quad (21)$$

Coupling with the local anisotropy is obtained by the substitution of η_{jl} instead of κ_{jl} . For small v both couplings are weakened by a factor v^2 . A typical transverse fluctuation (cf. Eq. (14)) contains η^2 in the numerator and $\sqrt{\kappa}$ in the denominator, so that the fluctuation is proportional to v^3 . Growth of transverse fluctuations at a decrease of the global anisotropy κ can be compensated by a choice of sufficiently small v so that the transverse fluctuations remain small and region of critical fluctuations is not entered. The global anisotropy is a convenient parameter for theoretical analysis. In particular, it makes expressions for transverse fluctuations finite. The analogy between κ and $\tau = (T - T_c)/T_c$ makes possible to use the theory of critical phenomena as a guidance. Unfortunately, in practice the anisotropy (deformation) of aerogel is difficult to control or to vary continuously. It is particularly difficult in a region of small deformation $\gamma = \Delta l/l \sim 10^{-2} \div 10^{-3}$, which is of interest. An uncontrolled deformation of such order could be present in the most of the experiments with ^3He in aerogel.

4 Discussion

Global anisotropy of aerogel lifts continuous degeneracy of superfluid ^3He . Direct manifestation of the anisotropy is orientation of the orbital part of the average order parameter. Another important effect is a suppression of transverse fluctuations of the order parameter, which otherwise are critical. There remain a basic question about the structure of the A-like phase in the isotropic aerogel. Taking isotropic state as a limit of vanishing anisotropy helps to understand its nature.

Different possibilities for a structure of the A-like phase in the isotropic limit are discussed in the current literature. One of them is the LIM state. According to Ref.⁵ it has to form via a first order phase transition at a certain value of anisotropy

κ . By the order of magnitude this value coincides with the borderline anisotropy κ_c below which transverse fluctuations become critical.

Another possibility can be guessed by the analogy with the other continuously degenerate systems with a quenched random anisotropy. It is the formation of QLRO state¹¹. On approach to this state when global anisotropy tends to zero the average order parameter is presumably fading continuously. In both cases only orientation of the order parameter is involved. A form of the order parameter does not change.

In the present paper the third possibility is discussed. It consists in the change of a form of the order parameter which decreases its coupling with the random anisotropy. This adjustment makes possible to maintain a long-range order in a limit of vanishing anisotropy. Realization of this possibility in the A-like phase of superfluid ^3He depends on a landscape of the unperturbed free energy f_0 . If the landscape is favorable deviations of the order parameter from the ABM form can start at much higher anisotropy then the estimated critical value for transition in the LIM state or in a state with the QLRO.

No comparison of the expected properties of the proposed state with the existing experimental data was made here because of a possible ambiguity introduced in the data by an uncontrolled deformation of aerogel. One can remark only that neither of the data rules out the third possibility.

Investigation of the A-like phase of ^3He in aerogel with a possibility of tuning deformation of aerogel to a very low level is presently one of the most challenging problem in the field.

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